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# Leading asymmetries in two-hadron production in $e^+e^-$ annihilation at the $Z$ pole

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## Abstract

We present the leading unpolarized and single spin asymmetries in inclusive two-hadron production in electron-positron annihilation at the  $Z$  pole. The azimuthal dependence in the unpolarized differential cross section of almost back-to-back hadrons is a leading  $\cos(2\phi)$  asymmetry, which arises solely due to the intrinsic transverse momenta of the quarks. An extensive discussion on how to measure this asymmetry and the accompanying time-reversal odd fragmentation functions is given. A simple estimate indicates that the asymmetry could be of the order of a percent. © 1998 Elsevier Science B.V.

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Recently, we have presented the results of the complete tree-level calculation of inclusive two-hadron production in electron-positron annihilation via one photon up to subleading order in  $1/Q$  [1], where the scale  $Q$  is defined by the (timelike) photon momentum  $q$  (with  $Q^2 \equiv q^2$ ) and given by  $Q = \sqrt{s}$ . The quantity  $Q$  had to be much larger than characteristic hadronic scales, but – being interested in effects at subleading order – we considered energies only well below the threshold for the production of  $Z$  bosons.

In this article we extend those results to electron-positron annihilation into a  $Z$  boson, such that the results can be used to analyze LEP-I data. We will neglect contributions from photon exchange and  $\gamma$ - $Z$  interference terms, which are known to be numerically irrelevant on the  $Z$  pole. Only leading order  $(1/Q)^0$  effects are discussed, since for  $Q \gtrsim M_Z$  the power corrections of order  $1/Q$  are expected to be completely negligible. Furthermore, we will focus on tree level, i.e., order  $(\alpha_s)^0$ . A rich structure nevertheless arises when taking into account the intrinsic transverse momentum of the quarks and, possibly, polarization of the detected hadrons in the final state. By accounting for intrinsic transverse momentum effects we extend the results in the analysis of Chen *et al.* [2], where no azimuthal asymmetries arising from transverse momenta have been considered.

For details of the calculation and the formalism we refer to [1]. We shortly repeat the essentials. We consider the process  $e^- + e^+ \rightarrow \text{hadrons}$ , where the two leptons (with momentum  $l$  for the  $e^-$  and  $l'$  for the  $e^+$ , respectively) annihilate into a  $Z$  boson with momentum  $q = l + l'$ , which is timelike with  $q^2 \equiv Q^2$ . Denoting

the momentum of the two outgoing hadrons by  $P_h$  ( $h = 1, 2$ ) we use invariants  $z_h = 2P_h \cdot q/Q^2$ . We will consider the case where the two hadronic momenta  $P_1$  and  $P_2$  do not belong to the same jet (i.e.,  $P_1 \cdot P_2$  is of order  $Q^2$ ). In principle, the momenta can also be considered as the jet momenta themselves, but then effects due to intrinsic transverse momentum will be absent. We will treat the production of hadrons of which the spin states are characterized by a spin vector  $S_h$  ( $h = 1, 2$ ), satisfying  $P_h \cdot S_h = 0$  and  $-1 \leq S_h^2 \leq 0$ . In this way we can treat the case of unpolarized final states or final state hadrons with spin-0 and spin-1/2. In the present article we will disregard the polarization of hadron two (summation over spins). The final states which have to be identified and analyzed for the effects we discuss are simpler than the ones investigated by Artru and Collins [3], who proposed to measure azimuthal correlations in four-hadron production.

The cross section (including a factor 1/2 from averaging over incoming polarizations) for two-particle inclusive  $e^+e^-$  annihilation is given by

$$\frac{P_1^0 P_2^0 d\sigma^{(e^+e^-)}}{d^3P_1 d^3P_2} = \frac{\alpha_w^2}{4((Q^2 - M_Z^2)^2 + \Gamma_Z^2 M_Z^2) Q^2} L_{\mu\nu} \mathcal{W}^{\mu\nu}, \quad (1)$$

with  $\alpha_w = e^2/(16\pi \sin^2\theta_w \cos^2\theta_w)$  and the helicity-conserving lepton tensor (neglecting the lepton masses and polarization) is given by

$$L_{\mu\nu}(l, l') = (g_V^{l2} + g_A^{l2})[2l_\mu l'_\nu + 2l_\nu l'_\mu - Q^2 g_{\mu\nu}] - (2g_V^l g_A^l)2i \epsilon_{\mu\nu\rho\sigma} l^\rho l'^\sigma, \quad (2)$$

where  $g_V^l, g_A^l$  denote the vector and axial-vector couplings of the  $Z$  boson to the leptons, respectively, and the hadron tensor is given by

$$\mathcal{W}_{\mu\nu}(q; P_1 S_1; P_2 S_2) = \int \frac{d^3P_X}{(2\pi)^3 2P_X^0} \delta^4(q - P_X - P_1 - P_2) H_{\mu\nu}(P_X; P_1 S_1; P_2 S_2), \quad (3)$$

with

$$H_{\mu\nu}(P_X; P_1 S_1; P_2 S_2) = \langle 0 | J_\mu(0) | P_X; P_1 S_1; P_2 S_2 \rangle \langle P_X; P_1 S_1; P_2 S_2 | J_\nu(0) | 0 \rangle, \quad (4)$$

where a summation over spins of the unobserved *out*-state  $X$  is understood.

In order to expand the lepton and hadron tensors in terms of independent Lorentz structures, it is convenient to work with vectors orthogonal to  $q$ . A normalized timelike vector  $\hat{t}$  is defined by the boson momentum  $q$  and a normalized spacelike vector  $\hat{z}$  is defined by  $\tilde{P}^\mu = P^\mu - (P \cdot q/q^2) q^\mu$  for one of the outgoing momenta, say  $P_2$ ,

$$\hat{t}^\mu \equiv \frac{q^\mu}{Q}, \quad (5)$$

$$\hat{z}^\mu \equiv \frac{Q}{P_2 \cdot q} \tilde{P}_2^\mu = 2 \frac{P_2^\mu}{z_2 Q} - \frac{q^\mu}{Q}. \quad (6)$$

Vectors orthogonal to  $\hat{t}$  and  $\hat{z}$  are obtained with help of the tensors

$$g_\perp^{\mu\nu} \equiv g^{\mu\nu} - \hat{t}^\mu \hat{t}^\nu + \hat{z}^\mu \hat{z}^\nu, \quad (7)$$

$$\epsilon_\perp^{\mu\nu} \equiv -\epsilon^{\mu\nu\rho\sigma} \hat{t}_\rho \hat{z}_\sigma = \frac{1}{(P_2 \cdot q)} \epsilon^{\mu\nu\rho\sigma} P_{2\rho} q_\sigma. \quad (8)$$

Since we have chosen hadron two to define the longitudinal direction, the momentum of hadron one can be used



To leading order the expression for the hadron tensor, including quarks and antiquarks, is

$$\mathcal{W}^{\mu\nu} = 3 \int dp^- dk^+ d^2 p_T d^2 k_T \delta^2(\mathbf{p}_T + \mathbf{k}_T - \mathbf{q}_T) \text{Tr}(\overline{\Delta}(p) V^\mu \Delta(k) V^\nu) \Big|_{p^+ k^-} + \left( \begin{matrix} q \leftrightarrow -q \\ \mu \leftrightarrow \nu \end{matrix} \right), \quad (17)$$

where  $V^\mu = g_V \gamma^\mu + g_A \gamma_5 \gamma^\mu$  is the  $Z$  boson-quark vertex. We have omitted flavor indices and summation. The correlation functions  $\Delta$  and  $\overline{\Delta}$  are given by [4]:

$$\Delta_{ij}(k) = \sum_X \frac{1}{(2\pi)^4} \int d^4 x e^{ik \cdot x} \langle 0 | \psi_i(x) | P_1, S_1; X \rangle \langle P_1, S_1; X | \overline{\psi}_j(0) | 0 \rangle, \quad (18)$$

$$\overline{\Delta}_{ij}(p) = \sum_X \frac{1}{(2\pi)^4} \int d^4 x e^{-ip \cdot x} \langle 0 | \overline{\psi}_j(0) | P_2, S_2; X \rangle \langle P_2, S_2; X | \psi_i(x) | 0 \rangle \quad (19)$$

and the quark momentum  $k$  (and similarly for  $p$ ) and the polarization vector  $S_1$  (from now on we omit  $S_2$ ) are decomposed as follows:

$$k \approx \frac{1}{z} P_1 + k_T, \quad (20)$$

$$S_1 \approx \frac{\lambda_1}{M_1} P_1 + S_{1T}. \quad (21)$$

To leading order in  $1/Q$  one has that  $z = z_1$ . The (partly integrated) correlation function  $\Delta$  is parametrized as:

$$\frac{1}{z} \int dk^+ \Delta(P_1, S_1; k) \Big|_{k^- = P_1^-/z, \mathbf{k}_T} = \frac{M_1}{P_1^-} \left\{ D_1 \frac{\not{P}_1}{M_1} + D_{1T}^\perp \frac{\epsilon_{\mu\nu\rho\sigma} \gamma^\mu P_1^\nu k_T^\rho S_{1T}^\sigma}{M_1^2} - G_{1s} \frac{\not{P}_1 \gamma_5}{M_1} \right. \\ \left. - H_{1T} \frac{i \sigma_{\mu\nu} \gamma_5 S_{1T}^\mu P_1^\nu}{M_1} - H_{1s}^\perp \frac{i \sigma_{\mu\nu} \gamma_5 k_T^\mu P_1^\nu}{M_1^2} + H_{1s}^\perp \frac{\sigma_{\mu\nu} k_T^\mu P_1^\nu}{M_1^2} \right\}, \quad (22)$$

where the shorthand notation  $G_{1s}$  (and similarly for  $H_{1s}^\perp$ ) stands for the combination

$$G_{1s}(z, \mathbf{k}_T) = \lambda_1 G_{1L} + G_{1T} \frac{(\mathbf{k}_T \cdot \mathbf{S}_{1T})}{M_1}. \quad (23)$$

We parametrize the antiquark correlation function  $\overline{\Delta}$  in the same way, except that the distribution functions are overlined and the obvious replacements of momenta are done.

The functions  $D_1, \dots$  in Eq. (22) and  $G_{1L}, G_{1T}, \dots$  in  $G_{1s}, \dots$  are called fragmentation functions. One wants to express the fragmentation functions in terms of the hadron momentum, hence, the arguments of the fragmentation functions are chosen to be the lightcone (momentum) fraction  $z = P_1^-/k^-$  of the produced hadron with respect to the fragmenting quark and  $\mathbf{k}'_T \equiv -z\mathbf{k}_T$ , which is the transverse momentum of the hadron in a frame where the quark has no transverse momentum. In order to switch from quark to hadron transverse momentum a Lorentz transformation leaving  $k^-$  and  $P_1^-$  unchanged needs to be performed. The fragmentation functions are real and in fact, depend on  $z$  and  $\mathbf{k}'_T$  only.

We note that after integration over  $\mathbf{k}_T$  several functions disappear. In the case of  $\text{Tr}(\Delta i \sigma^{i-} \gamma_5)$  a specific combination remains, namely  $H_1 \equiv H_{1T} + (\mathbf{k}_T^2/2M_1^2) H_{1T}^\perp$ .

The choice of factors in the definition of fragmentation functions is such that  $\int dz d^2 \mathbf{k}'_T D_1(z, \mathbf{k}'_T) = N_h$ , where  $N_h$  is the number of produced hadrons.

Note that the decay probability for an unpolarized quark with non-zero transverse momentum can lead to a transverse polarization in the production of spin-1/2 particles. This polarization is orthogonal to the quark transverse momentum and the probability is given by the function  $D_{1T}^\perp$ . In the same way, oppositely

transversely polarized quarks with non-zero transverse momentum can produce unpolarized hadrons or spinless particles, with different probabilities. In other words: there can be a preference for one or the other transverse polarization direction of the quark (aligned or opposite relative to its transverse momentum) to fragment into an unpolarized hadron. This difference is described by the function  $H_1^\perp$ . It is the one appearing in the so-called Collins effect [5], which predicts a single transverse spin asymmetry in for instance semi-inclusive DIS, and arises due to intrinsic transverse momentum.

The functions  $D_{1T}^\perp$  and  $H_1^\perp$  are what are generally called ‘time-reversal odd’ functions. For a discussion on the meaning of this, we refer to Ref. [1] and earlier Refs. [6]; here we only remark that it does not signal a violation of time-reversal invariance of the theory, but rather the presence of final state interactions.

The cross sections are obtained from the hadron tensor after contraction with the lepton tensor

$$L^{\mu\nu} = (g_V^{l2} + g_A^{l2}) Q^2 \left[ - (1 - 2y + 2y^2) g_\perp^{\mu\nu} + 4y(1-y) \hat{z}^\mu \hat{z}^\nu - 4y(1-y) \left( \hat{l}_\perp^\mu \hat{l}_\perp^\nu + \frac{1}{2} g_\perp^{\mu\nu} \right) \right. \\ \left. - 2(1-2y) \sqrt{y(1-y)} \hat{z}^{[\mu} \hat{l}_\perp^{\nu]} \right] - (2g_V^l g_A^l) Q^2 \left[ + i(1-2y) \epsilon_\perp^{\mu\nu} - 2i \sqrt{y(1-y)} \hat{l}_\perp^\rho \epsilon_\perp^{\rho[\mu} \hat{z}^{\nu]} \right], \quad (24)$$

where  $\{\mu\nu\}$  indicates symmetrization of indices,  $[\mu\nu]$  indicates antisymmetrization. The fraction  $y$  is defined to be  $y = P_2 \cdot l / P_2 \cdot q \approx l^- / q^-$ , which in the lepton center of mass frame equals  $y = (1 + \cos\theta_2)/2$ , where  $\theta_2$  is the angle of hadron two with respect to the momentum of the incoming leptons.

Azimuthal angles inside the perpendicular plane are defined with respect to  $\hat{l}_\perp^\mu$ , defined to be the normalized perpendicular part of the lepton momentum  $l$ ,  $\hat{l}_\perp^\mu = l_\perp^\mu / (Q\sqrt{y(1-y)})$ :

$$\hat{l}_\perp \cdot a_\perp = -|a_\perp| \cos\phi_a, \quad (25)$$

$$\epsilon_\perp^{\mu\nu} \hat{l}_\perp^\mu a_{\perp\nu} = |a_\perp| \sin\phi_a, \quad (26)$$

for a generic vector  $a$ .

The vector and axial couplings to the  $Z$  boson are given by:

$$g_V^j = T_3^j - 2Q^j \sin^2\theta_w, \quad (27)$$

$$g_A^j = T_3^j, \quad (28)$$

where  $Q^j$  denotes the charge and  $T_3^j$  the weak isospin of particle  $j$  (i.e.,  $T_3^j = +1/2$  for  $j = u$  and  $T_3^j = -1/2$  for  $j = e^-, d, s$ ). Combinations of the couplings occurring frequently in the formulas are

$$c_1^j = (g_V^{j2} + g_A^{j2}), \\ c_2^j = (g_V^{j2} - g_A^{j2}), \quad j = \ell \text{ or } u, d, s \\ c_3^j = 2g_V^j g_A^j. \quad (29)$$

As well, we will use the following kinematical factors:

$$A(y) = \left( \frac{1}{2} - y + y^2 \right) \stackrel{cm}{=} (1 + \cos^2\theta_2)/4, \\ B(y) = y(1-y) \stackrel{cm}{=} \sin^2\theta_2/4, \\ C(y) = (1-2y) \stackrel{cm}{=} -\cos\theta_2. \quad (30)$$

We obtain in leading order in  $1/Q$  and  $\alpha_s$  the following expression for the cross section in case of unpolarized (or spinless) final state hadrons:

$$\frac{d\sigma(e^+e^- \rightarrow h_1 h_2 X)}{d\Omega dz_1 dz_2 d^2\mathbf{q}_T} = \sum_{a,\bar{a}} \frac{3\alpha_w^2 Q^2}{(Q^2 - M_Z^2)^2 + \Gamma_Z^2 M_Z^2} z_1^2 z_2^2 \left\{ \left( c_1' c_1^a A(y) - \frac{1}{2} c_3' c_3^a C(y) \right) \mathcal{F} \left[ D_1^a \bar{D}_1^a \right] \right. \\ \left. + \cos(2\phi_1) c_1' c_2^a B(y) \mathcal{F} \left[ \left( 2 \hat{\mathbf{h}} \cdot \mathbf{k}_T \hat{\mathbf{h}} \cdot \mathbf{p}_T - \mathbf{k}_T \cdot \mathbf{p}_T \right) \frac{H_1^{\perp a} \bar{H}_1^{\perp a}}{M_1 M_2} \right] \right\}, \quad (31)$$

where  $d\Omega = 2dy d\phi^l$  and  $\phi^l$  gives the orientation of  $\hat{l}_\perp^\mu$ . We use the convolution notation

$$\mathcal{F} \left[ D^a \bar{D}^a \right] \equiv \int d^2\mathbf{k}_T d^2\mathbf{p}_T \delta^2(\mathbf{p}_T + \mathbf{k}_T - \mathbf{q}_T) D^a(z_1, z_1^2 \mathbf{k}_T^2) \bar{D}^a(z_2, z_2^2 \mathbf{p}_T^2). \quad (32)$$

The angle  $\phi_1$  is the azimuthal angle of  $\hat{\mathbf{h}}$  (see Fig. 1). In order to deconvolute these expressions we can define weighted cross sections

$$\langle W \rangle_A = \int \frac{d\phi'}{2\pi} d^2\mathbf{q}_T W \frac{d\sigma(e^+e^- \rightarrow h_1 h_2 X)}{d\Omega dz_1 dz_2 d^2\mathbf{q}_T}, \quad (33)$$

where  $W = W(Q_T, \phi_1, \phi_2, \phi_{S_1}, \phi_{S_2})$ . The subscript  $A$  denotes the polarization in the final state for hadron one, with as possibilities unpolarized ( $O$ , including the case of summation over spin), longitudinally polarized ( $L$ ) or transversely polarized ( $T$ ). We postpone the discussion of the additional structures and information accessible by measuring the polarization of one of the final state hadrons to the end of this letter.

Even without determining polarization of a final state hadron a subtle test of our understanding of spin transfer mechanisms in perturbative QCD can be done. The information on the production of a transversely polarized quark-antiquark pair, which subsequently fragment into unpolarized (or spinless) hadrons with probabilities depending on the orientation of the (anti)quark's spin vector relative to its transverse momentum, is contained in the  $\cos(2\phi_1)$  azimuthal asymmetry<sup>1</sup>. To access this information we utilize the weighted cross sections

$$\langle 1 \rangle_O = \frac{3\alpha_w^2 Q^2}{(Q^2 - M_Z^2)^2 + \Gamma_Z^2 M_Z^2} \sum_{a,\bar{a}} \left( c_1' c_1^a A(y) - \frac{1}{2} c_3' c_3^a C(y) \right) D_1^a(z_1) \bar{D}_1^a(z_2), \quad (34)$$

$$\left\langle \frac{Q_T^2}{4M_1 M_2} \cos(2\phi_1) \right\rangle_O = \frac{3\alpha_w^2 Q^2}{(Q^2 - M_Z^2)^2 + \Gamma_Z^2 M_Z^2} \sum_{a,\bar{a}} c_1' c_2^a B(y) H_1^{\perp(1)a}(z_1) \bar{H}_1^{\perp(1)a}(z_2), \quad (35)$$

where the  $\mathbf{k}_T^2$ -moments for a generic fragmentation function  $F$  are defined by

$$F^{(n)}(z_i) = z_i^2 \int d^2\mathbf{k}_T \left( \frac{\mathbf{k}_T^2}{2M_i^2} \right)^n F(z_i, z_i^2 \mathbf{k}_T^2). \quad (36)$$

We now like to focus on the weighted cross section defined in Eq. (35) and discuss its possible measurement. In order to be able to observe the  $\cos(2\phi)$  dependence one must look at two jet events in unpolarized electron-positron scattering. In each jet one identifies a fast hadron with momentum fractions  $z_1$  and  $z_2$  respectively. One of the hadrons (say two) together with the leptons determines the lepton scattering plane as is

<sup>1</sup> This asymmetry is not to be confused with the  $\cos(2\phi)$  asymmetry found by Berger [7], which is  $1/Q^2$  suppressed.

indicated in Fig. 1. In the lepton CM system hadron two determines the  $\hat{z}$ -direction with respect to which the azimuthal angles are measured. One needs in particular the azimuthal angle  $\phi_1$  of the other hadron (one) as well as its transverse momentum  $\mathbf{P}_{1\perp}$ , which determines  $Q_T = |\mathbf{P}_{1\perp}|/z_1$ . The  $\cos(2\phi)$  angular dependence then can be analyzed by calculating the weighted cross section of Eq. (35).

For an order of magnitude estimate, we consider the situation of the produced hadrons being a  $\pi^+$  and a  $\pi^-$ . Furthermore, we assume  $D_1^{u \rightarrow \pi^+}(z) = D_1^{\bar{d} \rightarrow \pi^+}(z)$  ( $D_1^{d \rightarrow \pi^-}(z) = D_1^{\bar{u} \rightarrow \pi^-}(z)$ , respectively) and neglect unfavored fragmentation functions like  $D^{d \rightarrow \pi^+}(z)$  etc.; and similar for the time-reversal odd functions. The equalities for the  $D_1$  functions seem quite safe on grounds of isospin and charge conjugation, the same assumptions might be non-trivial for the  $H_1^\perp$  functions. As a consequence of these assumptions the fragmentation functions can be taken outside the flavor summation, and we obtain

$$\left\langle \frac{Q_T^2}{4M_1 M_2} \cos(2\phi_1) \right\rangle_o = F(y) \frac{H_1^{\perp(1)}(z_1)}{D_1(z_1)} \frac{H_1^{\perp(1)}(z_2)}{D_1(z_2)} \langle 1 \rangle_o, \quad (37)$$

where

$$F(y) = \frac{\sum_{a=u,\bar{d}} c_1^e c_2^a B(y)}{\sum_{a=u,\bar{d}} \left( c_1^e c_1^a A(y) - \frac{1}{2} c_3^e c_3^a C(y) \right)}. \quad (38)$$

This factor is shown in Fig. 2 as a function of the center of mass angle  $\theta_2$  (we use  $\sin^2\theta_w = 0.2315$  [8]). At an angle close to  $90^\circ$  we observe the largest effect. In order to get an estimate of the true asymmetry at the level of count rates, one should compare Eq. (35) with the weighted cross section  $\langle Q_T^2/4M_1 M_2 \rangle_o$ . To estimate the ratio of those two quantities, we use as argued in Ref. [9], for the ratio of the fragmentation functions  $H_1^{\perp(1)}(z_1)/D_1(z_1) = \mathcal{O}(1)$ , although this is likely an optimistic estimate. From the average transverse momentum squared of produced pions in one jet, for which we take  $0.5 \text{ (GeV/c)}^2$  [10], one obtains an estimate for the average transverse momentum of pions in jet one with respect to a given pion in jet two. This leads at  $z_1 = z_2 = 1/2$  to  $\langle Q_T^2/4M_\pi^2 \rangle_o \approx 50 \langle 1 \rangle_o$  and consequently to an estimate at the percent level for the ratio  $\langle (Q_T^2/4M_\pi^2) \cos(2\phi_1) \rangle_o / \langle Q_T^2/4M_\pi^2 \rangle_o$ . Such an azimuthal dependence in the unpolarized cross section, however, may be detectable in present-day electron-positron scattering experiments.

The situation where hadron two is taken to be a jet, which in this back-to-back jet situation is equivalent to analyzing the azimuthal structure of hadrons inside a jet, is obtained by considering  $\bar{D}_1^a(z_2) = \delta(1 - z_2)$  and  $\bar{H}_1^{\perp(1)a}(z_2) = 0$ . This gives the familiar result for  $\int dz_2 \langle 1 \rangle_o$  and it gives zero for the  $z_2$ -integrated  $\cos 2\phi$  azimuthal asymmetry.

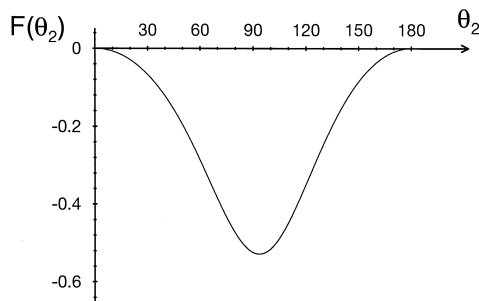


Fig. 2. Factor defined in Eq. (38) depending on the center of mass angle  $\theta_2$ .



Table 1

Weighted cross sections for  $S_1 \neq 0, S_2 = 0$ 

$W$	$A$	$\langle W \rangle_A \cdot \left[ 3 \alpha_w^2 Q^2 / \left( (Q^2 - M_Z^2)^2 + \Gamma_Z^2 M_Z^2 \right) \right]^{-1}$
1	$L$	$-\lambda_1 \sum_{a,\bar{a}} (c_1' c_3^a A(y) - \frac{1}{2} c_3' c_1^a C(y)) G_{1L}^a(z_1) \bar{D}_1^a(z_2)$
$(Q_T^2 / 4 M_1 M_2) \sin(2\phi_1)$	$L$	$\lambda_1 \sum_{a,\bar{a}} c_1' c_2^a B(y) H_{1L}^{\perp(1)a}(z_1) \bar{H}_1^{\perp(1)a}(z_2)$
$(Q_T / M_1) \sin(\phi_1 - \phi_{S_1})$	$T$	$ S_{1T}  \sum_{a,\bar{a}} (c_1' c_1^a A(y) - \frac{1}{2} c_3' c_3^a C(y)) D_{1T}^{\perp(1)a}(z_1) \bar{D}_1^a(z_2)$
$(Q_T / M_2) \sin(\phi_1 + \phi_{S_1})$	$T$	$- S_{1T}  \sum_{a,\bar{a}} c_1' c_2^a B(y) H_1^a(z_1) \bar{H}_1^{\perp(1)a}(z_2)$
$(Q_T^3 / 6 M_1^2 M_2) \sin(3\phi_1 - \phi_{S_1})$	$T$	$- S_{1T}  \sum_{a,\bar{a}} c_1' c_2^a B(y) H_{1T}^{\perp(2)a}(z_1) \bar{H}_1^{\perp(1)a}(z_2)$
$(Q_T / M_1) \cos(\phi_1 - \phi_{S_1})$	$T$	$ S_{1T}  \sum_{a,\bar{a}} (c_1' c_3^a A(y) - \frac{1}{2} c_3' c_1^a C(y)) G_{1T}^a(z_1) \bar{D}_1^a(z_2)$

The experimental determination of the polarization of (one of) the final state hadrons offers further opportunities to reveal the hadronic structure in terms of spin-dependent fragmentation functions. We assume in the following that the spin vector of hadron one, i.e.  $S_1$ , is known (reconstructed), having in mind the example of a produced  $\Lambda$  and its self-analyzing properties. We observe a rich structure of angular dependences due to polarization.

Again, weighted cross sections are the appropriate means to separate out specific functions. For instance, the weighted cross section

$$\left\langle \frac{Q_T}{M_2} \sin(\phi_1 + \phi_{S_1}) \right\rangle_T = |S_{1T}| \frac{3 \alpha_w^2 Q^2}{(Q^2 - M_Z^2)^2 + \Gamma_Z^2 M_Z^2} \sum_{a,\bar{a}} c_1' c_2^a B(y) H_1^a(z_1) \bar{H}_1^{\perp(1)a}(z_2) \quad (39)$$

picks out the term which is the closest analogue to the Collins effect [5] in semi-inclusive lepton-hadron scattering [11]. We note that a confirmation of the  $\cos(2\phi)$  asymmetry, also implies a confirmation of the Collins effect. A complete list of weighted cross sections at leading order is given in Table 1.

In conclusion, we have presented the leading asymmetries in inclusive two-hadron production in electron-positron annihilation at the  $Z$  pole. We have investigated unpolarized and single spin asymmetries. We included the effects of intrinsic transverse momentum and in this sense our results are an extension of those of Ref. [2]. The azimuthal dependence in the unpolarized differential cross section is a  $\cos(2\phi)$  asymmetry, which arises solely due to the intrinsic transverse momenta of the quarks. An extensive discussion on how to measure this asymmetry and the accompanying time-reversal odd fragmentation functions is given. A simple estimate indicates that the asymmetry could be at the percent level, hence it can perhaps be observed in present-day electron-positron scattering experiments. In confirming the existence of this asymmetry one also confirms the Collins effect, without the need of a polarization measurement.

*Note added:* In a preliminary study [12] a similar correlation in back-to-back jets was already experimentally investigated. We find that it involves moments of the functions  $H_1^\perp$  and  $\bar{H}_1^\perp$ , different from the ones in our correlation. In this study no significant result was found using the 1991 to 1994 LEP data. In this analysis *three* momenta in the final state need to be determined, namely besides two hadron momenta also the jet axis, and hence there are two azimuthal angles,  $\phi$  and  $\phi'$ , yielding a  $\cos(\phi + \phi')$  asymmetry.

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